

Exclusive Scattering at ELFE ¹

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RESUME:

Il était une fois dans le royaume *Standard* un preux chevalier rêvant de conquérir une mystérieuse province appelée *confinement*. Sa quête durait depuis plusieurs dizaines d'années. Un jour, un mystérieux *elfe* vint lui proposer son aide...C'est le premier chapitre de ses aventures que nous ouvrons aujourd'hui, dans ce livre situé aux *frontières* de cette mystérieuse province.

ABSTRACT:

The theoretical framework of hard exclusive reactions is reviewed with special emphasis on the Elfe project program. Perturbative QCD studies have shown that factorization properties allow to separate well-defined non perturbative objects which are crucial in the understanding of confinement dynamics from perturbatively calculable hard processes. The applicability of this factorization in a definite energy domain is controlled by some definite statements, as the dimensional counting rules, the helicity conservation law and the appearance of color transparency. The few data available indicate that the Elfe parameters indeed correspond to this well defined physics.

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Fig.1: The proton magnetic form factor

1 Introduction

The theory of hard elastic scattering in Quantum Chromodynamics (QCD) has evolved considerably over many years of work. Currently there exists a self-consistent perturbative description, with a specific factorization method for separating the hard scattering from non-perturbative wave functions. A well-known procedure using the “quark-counting” diagrams [1] has been given by LePage and Brodsky [2]. The value of Q^2 at which the scaling limit appears is matter of debate and of detailed phenomenological analysis.

The magnetic proton form factor (Fig.1) is the best known exclusive observable. Scaling obviously settles at Q^2 around 10 GeV^2 . Note that Fig.1 has a linear scale; the slight decrease at large Q^2 is understood as coming from logarithmic corrections which are well under control (see below). The analysis of the pion form factor is subject to more model dependent analysis. It however seems to scale at even lower values of Q^2 .

In the case of the Deuteron, Brodsky and Chertok [3] propose a very early scaling. A recent analysis[4] shows however that if this is the case then the deuteron distribution amplitude is very different from a proton-neutron bound state, for instance that it contains, and even is

dominated by, non nucleonic components. In the case of exclusive reactions at fixed angle, for instance real or virtual Compton Scattering or meson photo- and electro-production, the same arguments predict power behaved cross sections following QCD improved counting rules. The measure of the Q^2 or s - dependence of as many as possible hard exclusive observables at high energy is thus of utmost importance for the understanding of the applicability of the QCD framework at accessible energies.

A supplementary tool to disentangle the high transfer regime from the low energy domain is the predicted occurrence of color transparency when reactions on nuclei are performed. There is still no decisive experimental evidence of this phenomenon although the admittedly controversial explanation [5] of pp data at BNL [6] and recent data on diffractive heavy vector meson production at Fermilab [7] are certainly indicative.

2 Factorization : The example of form factors

The spacelike form factor measures the ability of a pion to absorb a virtual photon (carrying a momentum q with $q^2 = -Q^2 < 0$) while remaining intact. It is defined by the formula:

$$\langle \pi(p') | J^\mu | \pi(p) \rangle = e_\pi \cdot F(Q^2) \cdot (p + p')^\mu, \quad (1)$$

where e_π is the pion electric charge.

In the hard scattering regime, that is when Q is very high with respect to the low energy scales of the theory (the QCD scale Λ and the pion mass), Brodsky and Lepage have motivated the following three step picture for the process [2]:

- the pion exhibits a valence quark-antiquark “soft” (see below) state,
- which interacts with the hard photon leading to another soft state,
- which forms the final pion.

This leads to the convolution formula:

$$F(Q^2) = \psi_{in} * T_H * \psi_{out}^* \quad (2)$$

and the graphical representation of Fig.2.

Fig.2: The factorization of a hard scattering amplitude

The most important feature of this picture is that it separates hard from soft dynamics. The amplitude T_H , the *interaction*, reflects the hard transformation due to the absorption of the photon and is hopefully calculable in perturbative QCD, because the effective couplings are small in this regime due to the asymptotic freedom. The amplitude ψ , the *wave function*, which depends on low energy dynamics is outside of the domain of applicability of perturbative QCD and is, at present, far from being fully understood from the theory. It is however process independent and contains much information on confinement dynamics. Factorization proofs legitimate this picture [2].

2.1 Sudakov effects

The need of a careful factorization is due to the infrared behavior of QCD: technically, large logarithms ($\sim \ln(Q/\lambda)$) appear in the renormalized one-loop corrections to naive “tree-graph” (λ is some infrared cut off needed to regularize soft and/or collinear divergences). As in the renormalization procedure, if factorization holds, these large corrections should be absorbed, here in the redefinition of the wave function. The proof of factorization and its consequences upon the wave functions are studied in the pattern of the renormalization group. Without entering into a detailed discussion, let us sketch the procedure (see [8] for more on this leading logarithms calculation and also for the renormalization group treatment).

Fig.3: The tree graph for the meson form factor

The first step is to compute the naive hard amplitude, that is consider the tree graph of Fig.3 , and the three other graphs related to it by C and T symmetries. One finds, with notations explained on Fig.3:

$$T_H = 16\pi\alpha_S C_F \frac{xQ^2}{xQ^2 + \mathbf{k}^2 - i\varepsilon} \frac{1}{xyQ^2 + (\mathbf{k} - \mathbf{l})^2 - i\varepsilon}, \quad (3)$$

where all quark momentum components are kept. Note that we have done the usual projection onto the pion S wave state: $\psi_\pi(p) \propto \frac{1}{\sqrt{2}}\gamma^5 \not{p}$ and used the C symmetry of the wave function. $C_F = 4/3$ is the color factor, while α_S is the QCD effective coupling at the renormalization point μ .

Let us now examine one loop corrections to T_H In axial gauge, it turns out that the relevant graphs to consider are those of Fig.4.: they directly lead to the wave function correction, in the “double logarithms” or Sudakov region (namely: $\lambda \ll |\mathbf{q}| \ll u\frac{Q}{\sqrt{2}} \ll x\frac{Q}{\sqrt{2}}$, u and \mathbf{q} being respectively the light-cone fraction and transverse gluon momentum relatively to the pion):

$$\begin{aligned} \psi^{(1)}(x, \mathbf{k}) = & \frac{C_F}{2\pi^2} \int_\lambda^{xQ/\sqrt{2}} \frac{d^2\mathbf{q}}{\mathbf{q}^2} \alpha_S(\mathbf{q}^2) \int_{|\mathbf{q}|\sqrt{2}/Q}^x \frac{du}{u} \{ \psi^{(0)}(x-u, \mathbf{k} + \mathbf{q}) - \psi^{(0)}(x, \mathbf{k}) \} \\ & + \frac{C_F}{2\pi^2} \int_\lambda^{\bar{x}Q/\sqrt{2}} \frac{d^2\mathbf{q}}{\mathbf{q}^2} \alpha_S(\mathbf{q}^2) \int_{|\mathbf{q}|\sqrt{2}/Q}^{\bar{x}} \frac{du}{u} \{ \psi^{(0)}(x+u, \mathbf{k} + \mathbf{q}) - \psi^{(0)}(x, \mathbf{k}) \}, \end{aligned} \quad (4)$$

where $\bar{x} = 1 - x$ and the first term in the difference comes from vertex-like corrections and

Fig.4: One loop graphs for the meson form factor

the second one from self energy ones; in the infrared region some partial cancellations occur between these corrections, but the cancellation is not complete.

To pursue this analysis, it is convenient to define the Fourier transform in the transverse plane:

$$\hat{\psi}(x, \mathbf{b}) = \int d^2\mathbf{k} e^{i\mathbf{k}\mathbf{b}} \psi(x, \mathbf{k}), \quad (5)$$

and to separate transverse and longitudinal variations of the wave function. One finds, omitting for the moment the second term in Eq. (4):

$$\begin{aligned} \hat{\psi}^{(1)}(x, \mathbf{b}) &= \frac{C_F}{2\pi^2} \left(\int \frac{d^2\mathbf{q}}{\mathbf{q}^2} \alpha_S(\mathbf{q}^2) (e^{-i\mathbf{q}\mathbf{b}} - 1) \int \frac{du}{u} \right) \hat{\psi}^{(0)}(x, \mathbf{b}) \\ &+ \frac{C_F}{2\pi^2} \int \frac{d^2\mathbf{q}}{\mathbf{q}^2} \alpha_S(\mathbf{q}^2) e^{-i\mathbf{q}\mathbf{b}} \int \frac{du}{u} \left(\hat{\psi}^{(0)}(x - u, \mathbf{b}) - \hat{\psi}^{(0)}(x, \mathbf{b}) \right). \end{aligned} \quad (6)$$

This equation contains the typical corrections one has to consider in a hard process when dealing with either a big ($\gg 1/Q$) or a small ($< 1/Q$) neutral object.

The transverse behavior at large distance is driven by the first term of the previous equation, thanks to the vanishing of the summation with the oscillating components. This occurs when $b = |\mathbf{b}|$ is greater than a few times the inverse of the upper bound of the corresponding integral: $xQ/\sqrt{2}$. As a consequence, in the remaining expression, the infrared cut-off λ can be replaced by the natural one $1/b$, above which the vertex and self energy corrections do not

compensate one another. Thus we get [9]:

$$\hat{\psi}^{(1)} = -s(x, Q, b) \hat{\psi}^{(0)}, \quad s = \frac{C_F}{2\beta} \ln \frac{xQ}{\sqrt{2}} \left(\ln \frac{\ln xQ/\sqrt{2}}{\ln 1/b} - 1 + \frac{\ln 1/b}{\ln xQ/\sqrt{2}} \right), \quad (7)$$

with $\beta = (11 - \frac{2n_f}{3})/4$, n_f being the number of quark flavors. Here and in the following, it is understood that the energies and inverse separations are in units of the natural QCD scale Λ_{QCD} . After the resummation of the ladder structure to all orders, the above Sudakov factor exponentiates. Taking into account the term obtained with the substitution $x \rightarrow \bar{x} \equiv 1 - x$, we get:

$$\hat{\psi}(x, b, Q) = e^{-s(x, b, Q) - s(\bar{x}, b, Q)} \hat{\psi}^{(0)}(x, b, Q). \quad (8)$$

Thus we get a strong suppression of the effective wave function as $b \rightarrow 1/\Lambda$, whatever the fraction x is, provided that Q is reasonably large. **The remaining object $\hat{\psi}^{(0)}$ is a soft component to start with.** It is soft in the sense that it does not include loop-corrections harder than $1/b$. One may modelize it by including some b behavior or simply relate it to the distribution amplitude [8] setting:

$$\hat{\psi}^{(0)}(x, b) \approx \int_0^{1/b} \psi(x, \mathbf{k}) d\mathbf{k} = \varphi(x; 1/b). \quad (9)$$

2.2 Leading log analysis

The first term in Eq.6 is negligible when the oscillating term remains close to 1 in the range of integration. This happens for b a few times less than $\max^{-1}(xQ, \bar{x}Q)$. In this case, soft divergences cancel one another and one finds:

$$\hat{\psi}^{(1)}(x) = \xi \frac{C_F}{2} \int_0^1 dx' \left\{ \frac{\hat{\psi}^{(0)}(x') - \hat{\psi}^{(0)}(x)}{x - x'} \theta(x - x') + \frac{\hat{\psi}^{(0)}(x') - \hat{\psi}^{(0)}(x)}{x' - x} \theta(x' - x) \right\}, \quad (10)$$

with the notation:

$$\xi = \frac{1}{\pi^2} \int_\lambda^Q \frac{d^2 \mathbf{q}}{\mathbf{q}^2} \alpha_S(\mathbf{q}^2) \sim \frac{1}{\beta} \ln \left(\frac{\ln Q}{\ln \lambda} \right). \quad (11)$$

We displayed this equation in a slightly different form than in the large b case to explicitly show that equation 4, in the limit of small b , is related to the distribution evolution proposed in [2]. Let us shortly review how this comes[10]. The earlier simpler factorization formula for exclusive processes [2] is easily derived from the previous treatment if one assumes that neither the wave function nor the hard amplitude give important contribution to the form factor when

the transverse momenta are big. Neglecting all transverse momenta in T_H leads therefore to consider the k_T -integrated quantity:

$$\varphi(x) = \int d^2\mathbf{k} \psi(x, \mathbf{k}) \quad (12)$$

This distribution amplitude related to the wave function at $b = 0$ has a dependence in Q associated with the remaining collinear divergences in Eq.6. Indeed, the exponentiated form of this convolution equation, once it is written for the distribution φ and generalized to other regions than the Sudakov one, leads to the celebrated expansion of $\varphi(x, Q)/x\bar{x}$ in a linear combination of running logarithms together with Gegenbauer polynomials.

Before turning to this point, let us examine a useful constraint on the distribution amplitude. In fact we know something about the wave function, and this is because the pion electroweak decay is a *very* short distance processus, which requires the pion to be in its valence state. The pion decay is described by the matrix element:

$$\langle 0 | \bar{q}_d(0) \gamma^\mu (1 - \gamma^5) q_u(0) | \pi^+(p) \rangle = f_\pi p^\mu, \quad (13)$$

where f_π has been measured as 133 MeV.

The amplitude at zero distance may be written as:

$$\langle 0 | T(q_{u\alpha i}(0) \bar{q}_{d\beta j}(0)) | \pi^+(p) \rangle = \int_0^1 dx \frac{Q}{2\sqrt{2}\pi} \int \frac{dk^- d\mathbf{k}_\perp}{(2\pi)^3} X(k), \quad (14)$$

which must be projected onto the tensor $\gamma^\mu(1 - \gamma^5)|_{\beta\alpha}\delta_{ji}$, to yield:

$$- \langle 0 | T(\bar{q}_{di}(0) \gamma^\mu (1 - \gamma^5) q_{ui}(0)) | \pi^+(p) \rangle = \int_0^1 \varphi(x) Tr \left(\frac{\gamma^5 \not{p}}{4} \gamma^\mu (1 - \gamma^5) \right) \frac{\delta_{ij}}{3} \delta_{ji}. \quad (15)$$

We get:

$$\int_0^1 dx \varphi(x) p^\mu = f_\pi p^\mu, \quad (16)$$

so that the normalisation of the distribution amplitude is fixed.

Let us now see how leading logarithms are resummed into the distribution amplitude, in a way which is much reminiscent of the now standard Altarelli-Parisi evolution equations for

structure functions in deep inelastic reactions. In axial gauge, one may isolate big logarithmic factors in ladder-type diagrams, leading to an expansion:

$$\begin{aligned}\varphi_{\text{LL}}(x, Q) &= \varphi_0(x) + \kappa \int_0^1 du V_{q\bar{q} \rightarrow q\bar{q}}(u, x) \varphi_0(u) \\ &+ \frac{\kappa^2}{2!} \int_0^1 du V_{q\bar{q} \rightarrow q\bar{q}}(u, x) \int_0^1 du' V_{q\bar{q} \rightarrow q\bar{q}}(u', u) \varphi_0(u') + \dots\end{aligned}\quad (17)$$

where

$$\kappa = \frac{2}{\beta} \ln \frac{\alpha_S(\mu^2)}{\alpha_S(Q^2)},$$

and $V_{q\bar{q} \rightarrow q\bar{q}}$ is the following kernel

$$V_{q\bar{q} \rightarrow q\bar{q}}(u, x) = \frac{2}{3} \left\{ \frac{\bar{u}}{\bar{x}} \left(1 + \frac{1}{u-x} \right)_+ \theta(u-x) + \frac{u}{x} \left(1 + \frac{1}{x-u} \right)_+ \theta(x-u) \right\}, \quad (18)$$

where the $(\cdot)_+$ distribution comes from the colour neutrality of the meson.

The equation on φ may be rewritten as

$$\left(\frac{\partial \varphi}{\partial \kappa} \right)_x = \int_0^1 du V(u, x) \varphi(x, Q), \quad (19)$$

the solution of which has been known for a long time as:

$$\varphi(x, Q) = x(1-x) \sum_n \phi_n(Q) C_n^{(3/2)}(2x-1); \quad (20)$$

where the $C_n^{(m)}$ are Gegenbauer polynomials which verify:

$$\int_0^1 du u(1-u) V(u, x) C_n^{(3/2)}(2u-1) = A_n x(1-x) C_n^{(3/2)}(2x-1), \quad (21)$$

with eigenvalues A_n . One thus gets

$$\phi_n(Q) = \phi(\mu) e^{A_n \kappa} = \phi(\mu) \left(\frac{\alpha_S(\mu^2)}{\alpha_S(Q^2)} \right)^{2A_n/\beta}, \quad (22)$$

where exponents monotonously decrease, beginning with: $\frac{2A_0}{\beta} = 0$, $\frac{2A_2}{\beta} = -0,62$, ...

Using the previously derived normalization:

$$\int_0^1 dx \varphi(x, Q) = \phi_0(Q) \int_0^1 dx x(1-x) = \frac{\phi_0}{6} = f_\pi, \quad (23)$$

the expansion may be written as:

$$\varphi(x, Q) = 6f_\pi x(1-x) + \Phi_2(\ln Q^2)^{-0.62} x(1-x)(2x-1)^2 + \dots \quad (24)$$

We thus know the asymptotic pion distribution amplitude:

$$\varphi(x, Q) \sim_{Q \rightarrow \infty} 6f_\pi x(1-x). \quad (25)$$

We however do not know the realistic distribution amplitude at accessible energies, since the constants $\Phi_2, \dots \Phi_n$ cannot be derived from perturbative QCD. This is where the QCD sum rule approach enters [11] to yield values for moments of the distribution

$$\int_0^1 dx (2x-1)^{2n} \varphi(x, \mu), \dots \quad (26)$$

and then reconstruct the constants $\Phi_2, \dots \Phi_n$ [12].

3 Exclusive reactions at ELFE

Measuring the pion form factor, and more generally the dominant form factors for hadrons (F_K, G_N, \dots) just allows to get one constraint on the distribution amplitude of these hadrons. To go further, one needs to investigate scattering amplitudes at large angle which map out the x_i dependence of the wave functions. Elastic real and virtual Compton scattering ($e + p \rightarrow e + p + \gamma$) at high momentum transfer meet this goal since they only depend on the valence proton distribution amplitude (eventually Sudakov improved) analysed by a calculable hard amplitude. Different meson photo- and electro-production processes depend on more hadronic wave functions and are as thus probes of various ways quarks are confined in hadrons.

Subdominant form factors (such as F_2, \dots) and helicity violating processes (such as non-diagonal helicity matrix elements of produced vector mesons) are currently not understood with the same degree of rigor. Some progress has however been made [13] in a related process where it has been proposed that these observables measure a non-zero orbital angular momentum part of the valence wave-function.

3.1 Compton Scattering

To lowest order in the fine structure constant $\alpha \sim 1/137$, Virtual Compton Scattering (VCS) is described by the coherent sum of the amplitudes shown in Fig.5, namely a Bethe Heitler process (Fig.5b) where the final photon is radiated off the electron and a genuine VCS process (Fig.5a). Since the BH amplitude is calculable, at least once the elastic form factors

Fig.5: Amplitudes for virtual Compton scattering

$G_{Mp}(-t)$ and $G_{Ep}(-t)$ are known, its interference with VCS is a new source of information, not present in either real compton scattering or in electroproduction experiments (e.g. $(e, e'p)\pi^0$). The VCS amplitude depends upon three independent invariants, the usual choices being Q^2, s, t or Q^2, s, θ_{CM} .

Fig.6 displays the world's supply of high energy real Compton data on the proton, for $-t > 1\text{GeV}^2$. Although there are many experiments at high energy and low t , there is only the experiment of Shupe *et al.*, at large t . [14], The data are plotted as $s^6 d\sigma/dt$ vs. $\cos \theta_{CM}$ to illustrate the approach to the asymptotic scaling law. The most stringent test of the scaling law is obtained at $\theta_{CM} = 90^\circ$. Fitting the data to a $s^{-\alpha}$ power law results in $\alpha = 7.0 \pm 0.4$: a 2.5σ deviation from the $\alpha = 6.0$ prediction.

The perturbative QCD calculations of high energy real [15, 16] and virtual compton scattering [17] on the proton keep only the lowest order Feynman diagrams (there are 336 non-vanishing topologically distinct ways to couple two photons to three quarks with the exchange of two gluons and 42 diagrams with the three gluon vertex whose color factor however vanishes). Each quark entering or leaving the hard scattering carries a fraction x of the momentum of its parent proton and components of momentum along the three other directions. As long as x -fractions are fixed and non-zero, and at large enough momentum transfer, it is legitimate to further neglect, in the hard amplitude, components of each quark internal momenta which do

Fig.6: Real Compton Scattering at large angle

not lie along its parent proton , one thus gets

$$A = \phi_{(uud)} \otimes T_H(\{x\}, \{y\}) \otimes \phi'_{(uud)}(1 + O(M^2/t)), \quad (27)$$

3.2 A strategy for data analysis

A first way to analyse data is to compare experimental points to a calculation with a given distribution amplitude that any prejudiced theorist convinced you to choose. This is the way that the pQCD calculations by Kronfeld and Nizic [16] of real Compton scattering are shown in Fig.6 for various choices of the distribution amplitude. One observes the good sensitivity of the Compton cross section to the non perturbative quantity $\phi_{(uud)}$ we are primarily interested in.

A more model-independent way is to try to sort out the wave function directly from the data. Let us outline a possible strategy on the example of real Compton scattering. We first write the proton valence wave-function as the series derived from the leading logarithmic analysis, similar to Eq.24 for the pion case, that is in terms of Appel polynomials [12], as

$$\phi(x_i, Q) = 120x_1x_2x_3 \left\{ 1 + \frac{21}{2} \left(\frac{\alpha_S(Q^2)}{\alpha_S(Q_0^2)} \right)^{\lambda_1} A_1 P_1(x_i) + \frac{7}{2} \left(\frac{\alpha_S(Q^2)}{\alpha_S(Q_0^2)} \right)^{\lambda_2} A_2 P_2(x_i) + \dots \right\}, \quad (28)$$

where the slow Q^2 evolution entirely comes from renormalization group factors $\alpha_S(Q^2)^\lambda$, λ_i being calculated increasing numbers:

$$\lambda_1 = \frac{20}{9\beta} \quad , \quad \lambda_2 = \frac{24}{9\beta},$$

$P_i(x_j)$ are tabulated Apple polynomials

$$P_1(x_i) = x_1 - x_3 \quad , \quad P_2(x_i) = 1 - 3x_2, \dots$$

and A_i are unknown constants which measure the projection of the wave-function on the Apple polynomials:

$$A_i = \int_0^1 dx_1 dx_2 dx_3 \delta(x_1 + x_2 + x_3 - 1) \phi(x_i) P_i(x_i) \quad (29)$$

We then write the Compton differential cross-section as a sum of terms

$$A_i T_H^{ij}(\theta) A_j$$

where T_H^{ij} are integrals of the hard amplitude at some given scattering angle θ multiplied by the two Apple polynomials $A_i(x)$ and $A_j(y)$ over light-cone variables x and y , which have a rather awful analytical expression but can easily be electronically stored.

Sorting out the valence wave function of the proton from the data amounts then to determine through a maximum of likelihood method the parameters A_i restricting to something like ten terms in the expansion of Eq.28. A direct test of the validity of the approach is then to explore both real and virtual Compton scattering data which should be understood with the same series of A_i 's.

3.3 Other processes

Photo- and electro-production of mesons at large angle will enable to probe π and ρ distribution amplitudes in much the same way. The production of $K\Lambda$ final states will enable to enter strange quark production, thus selecting few diagrams in the hard process. Not much theoretical analysis of these possibilities has however been worked out except under the simplifying assumptions of the diquark model [18].

4 Color Transparency

4.1 The idea

The concept of Color transparency [19, 20] has recently attracted much attention. This phenomenon illustrates the power of exclusive reactions to isolate simple elementary quark configurations. The experimental technique to probe these configurations is the following:

For a hard exclusive reaction, say electron scattering from a proton, the scattering amplitude at large momentum transfer Q^2 is suppressed by powers of Q^2 if the proton contains more than the minimal number of constituents. This is derived from the QCD based quark counting rules, which result from the factorization of wave-function-like distribution amplitudes. Thus protons containing only valence quarks participate in the scattering. Moreover, each quark, connected to another one by a hard gluon exchange carrying momentum of order Q , should be found within a distance of order $1/Q$. Thus, at large Q^2 one selects a very special quark configuration: all connected quarks are close together, forming a small size color neutral configuration sometimes referred to as a *mini hadron*. This mini hadron is not a stationary state and evolves to build up a normal hadron.

Such a color singlet system cannot emit or absorb soft gluons which carry energy or momentum smaller than Q . This is because gluon radiation — like photon radiation in QED — is a coherent process and there is thus destructive interference between gluon emission amplitudes by quarks with "opposite" color. Even without knowing exactly how exchanges of soft gluons and other constituents create strong interactions, we know that these interactions must be turned off for small color singlet objects.

An exclusive hard reaction will thus probe the structure of a *mini hadron*, i.e. the short distance part of a minimal Fock state component in the hadron wave function. This is of primordial interest for the understanding of the difficult physics of confinement. First, selecting the simplest Fock state amounts to the study of the confining forces in a colorless object in the "quenched approximation" where quark-antiquark pair creation from the vacuum is forbidden. Secondly, letting the mini-state evolve during its travel through different nuclei of various sizes allows an indirect but unique way to test how the squeezed mini-state goes back to its full size and complexity, i.e. how quarks inside the proton rearrange themselves spatially to

”reconstruct” a normal size hadron. In this respect the observation of baryonic resonance production as well as detailed spin studies are mandatory.

To the extent that the electromagnetic form factors are understood as a function of Q^2 , $eA \rightarrow e'(A-1)p$ experiments will measure the color screening properties of QCD. The quantity to be measured is the transparency ratio T_r which is defined as:

$$T_r = \frac{\sigma_{Nucleus}}{Z\sigma_{Nucleon}} \quad (30)$$

At asymptotically large values of Q^2 , dimensional estimates suggest that T_r scales as a function of $A^{1/3}/Q^2$ [21]. The approach to the scaling behavior as well as the value of T_r as a function of the scaling variable determine the evolution from the pointlike configuration to the complete hadron. This highly interesting effect can be measured in an $e, e'p$ reaction that provides the best chance for a *quantitative* interpretation.

4.2 Available data

Experimental data on color transparency are very scarce but worth considering in detail. The first piece of evidence for something like color transparency came from the Brookhaven experiment on pp elastic scattering at 90 CM in a nuclear medium [6]. These data lead to a lively debate with no unanimous conclusion. The problem is that hadron hadron elastic scattering is not an as-well clear cut case of short distance process as the electromagnetic form factors discussed above. There are indeed infrared sensitive processes (the so-called independent scattering mechanism) which allow not so small protons to scatter elastically. The phenomenon of colour transparency is thus replaced by a *nuclear filtering* process: elastic scattering in a nucleus filters away the big component of the nucleon wave function and thus its contribution to the cross-section. Since the presence of these two competing processes had been analysed [22] as responsible of the oscillating pattern seen in the scaled cross-section $s^{10}d\sigma/dt$, an oscillating color transparency ratio emerges (see Fig.7)

The SLAC NE18[23] experiment recently measured the color transparency ratio up to $Q^2 = 7GeV^2$, without any observable increase. This casts doubts on the most optimistic views on very early dominance of point like configurations and emphasizes the importance of a sufficient boost to prevent small states to dress-up to quickly, then losing their ability to escape

Fig.7: Oscillating scaled cross-section R_1 (a)
and transparency ratio (b) for pp elastic scattering at 90 degrees[5].

freely the nucleus.

The diffractive electroproduction of heavy vector mesons at Fermilab [7] recently showed an interesting increase of the transparency ratio for data at $Q^2 = 7\text{GeV}^2$. In this case the boost is high since the lepton energy is around $E \simeq 200\text{GeV}$ but the problem is to disentangle diffractive from inelastic events.

5 Conclusion

Exclusive Scattering is certainly a central issue of the ELFE program and the beautiful theoretical progresses we have witnessed during these last few years eagerly await experimental data to be confronted with. We at last have a chance to pierce the secret of how quarks manage to build up hadrons by a thorough study of rare but illuminating processes. The new techniques now available for accelerating intense electron beams with a very high duty factor are quite fascinating. They will allow us to open a new chapter in the study of this challenging object we are all made of, namely the *proton*.

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